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Strong and Electromagnetic Decays of Two New Λ_c^* Baryons

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Abstract

Two recently discovered excited charm baryons are studied within the framework of Heavy Hadron Chiral Perturbation Theory. We interpret these new baryons which lie 308 MeV and 340 MeV above the Λ_c as $I = 0$ members of a P-wave spin doublet. Differential and total decay rates for their double pion transitions down to the Λ_c ground state are calculated. Estimates for their radiative decay rates are also discussed. We find that the experimentally determined characteristics of the Λ_c^* baryons may be simply understood in the effective theory.

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1. Introduction

The discovery of the first excited charm baryon has recently been announced by the ARGUS, CLEO and E687 groups [1–3]. The new state lies approximately 340 MeV above the Λ_c(2286 MeV) and decays to it via double pion emission. Although its spin, isospin and parity are not yet known, this new charmed baryon has been preliminarily interpreted as a Λ_c^* resonance. CLEO has further reported evidence for a second Λ_c^* excitation at 308 MeV above Λ_c [4]. The second resonance also decays through a double pion mode that is consistent with the two step process $\Lambda_c^* \to \Sigma_c \pi$ followed by $\Sigma_c \to \Lambda_c \pi$. In contrast, CLEO finds no evidence for an intermediate Σ_c in the decay of the first Λ_c^* excitation [2].

In this article, we will analyze these new baryon states and their dominant decay modes within the framework of Heavy Hadron Chiral Perturbation Theory (HHCPT). This hybrid effective theory represents a synthesis of Chiral Perturbation Theory and the Heavy Quark Effective Theory (HQET) and describes the low energy interactions between light Goldstone bosons and hadrons containing a heavy quark [5–9]. Since its development a few years ago, HHCPT has primarily been applied to the study of ground state charm and bottom hadrons. Ground state mesons and baryons are more tightly restricted by heavy quark spin symmetry than their excited counterparts. Moreover, experimental information has been much more sparse for the latter than the former. It is therefore not surprising that theorists have concentrated upon the lowest lying hadrons in the past. Now however that new data is being collected, it is worthwhile to broaden the scope of HHCPT and incorporate excited heavy hadrons into the effective theory.

The first excited heavy mesons and baryons are P-wave hadrons that carry one unit of orbital angular momentum. P-wave mesons have already been investigated within the HHCPT framework [10–12]. It is straightforward to extend the formalism and include P-wave baryons as well. A number of unknown couplings enter into the excited baryon sector which limits one's predictive power. But as we shall see, all the general characteristics of the two Λ_c^* baryons reported by ARGUS, CLEO and E687 are consistent with their being members of an excited spin symmetry doublet. Although our findings will be more qualitative than quantitative, we hope this work may help guide experimentalists as they continue to study these new charmed baryons.

Our paper is organized as follows. In section 2, we incorporate the lowest lying excited baryon doublet into the heavy baryon chiral Lagrangian. We then focus upon the two new Λ_c^* members of this doublet and analyze their strong interaction decays in section 3. Radiative transitions are discussed in section 4. Finally, we close in section 5 with some thoughts on future directions for investigation.

2. The Heavy Baryon Chiral Lagrangian

We begin by recalling some basic aspects of the baryon sector in Heavy Hadron Chiral Perturbation Theory [7,8]. Ground state baryons with quark content Qqq have zero orbital angular momentum and occur in two types depending upon the angular momentum j_ℓ of their light degrees of freedom. In the first case, the light brown muck is arranged in a symmetric $j_\ell = 1$ configuration which transforms as a sextet under flavor $SU(3)$. The spectators consequently couple with the heavy quark to form $J^P = \frac{1}{2}^+$ and $J^P = \frac{3}{2}^+$ S-wave bound states. When the heavy quark is taken to be charm, the spin-$\frac{1}{2}$ states are annihilated by velocity dependent Dirac operators $S^{ij}(v)$ whose individual components are given by

$$\begin{gathered} S^{11} = \Sigma_c^{++} \qquad S^{12} = \sqrt{\frac{1}{2}}\Sigma_c^{+} \qquad S^{22} = \Sigma_c^{0} \\ S^{13} = \sqrt{\frac{1}{2}}\Xi_c^{+\prime} \qquad S^{23} = \sqrt{\frac{1}{2}}\Xi_c^{0\prime} \\ S^{33} = \Omega_c^0. \end{gathered} \tag{2.1}$$

Their spin-$\frac{3}{2}$ counterparts are destroyed by corresponding $S^{*\,ij}_\mu(v)$ Rarita-Schwinger operators. In the second case, the light degrees of freedom form an antisymmetric $j_\ell = 0$ combination which transforms as a flavor antitriplet. Coupling with the heavy quark then yields $J^P = \frac{1}{2}^+$ baryons which we associate with the field $T_i(v)$. When $Q = c$, the individual components of T_i are the singly charmed baryons

$$T_1 = \Xi_c^0 \qquad T_2 = -\Xi_c^+ \qquad T_3 = \Lambda_c^+. \tag{2.2}$$

The complete spectrum of the first orbitally excited P-wave Qqq baryons is quite complicated. The lowest lying such hadrons correspond to bound states that have one unit of orbital angular momentum inserted between the heavy quark and light diquark pair. In this case, spin statistics constrain the light degrees of freedom to belong to either a $j_\ell = 1$ multiplet which transforms as a flavor antitriplet or else to $j_\ell = 0$, 1 or 2 multiplets which transform as flavor sextets. Nonrelativistic quark model calculations indicate that the antitriplet multiplet is isolated and lies significantly below all other P-wave states [13]. We will therefore only incorporate this lightest $j_\ell = 1$ multiplet into the chiral Lagrangian. We assign the Dirac and Rarita-Schwinger operators $R_i(v)$ and $R^*_{\mu i}(v)$ to its $J^P = \frac{1}{2}^-$ and $J^P = \frac{3}{2}^-$ states. As we shall see, the two newly discovered Λ_c^* baryons are well described as the $I = 0$ members of R and R^*_μ.

In the infinite heavy quark mass limit, it is useful to combine together the degenerate $J = \frac{1}{2}$ and $J = \frac{3}{2}$ members of the ground state sextet and excited antitriplet multiplets into the baryon "superfields" [8,14]

$$\begin{aligned} \mathcal{R}_{\mu i} &= \sqrt{\frac{1}{3}}(\gamma_\mu + v_\mu)\gamma^5 R_i + R^*_{\mu i} \\ \mathcal{S}^{ij}_\mu &= \sqrt{\frac{1}{3}}(\gamma_\mu + v_\mu)\gamma^5 S^{ij} + S^{*\,ij}_\mu . \end{aligned} \tag{2.3}$$

The $\mathcal{T}_i$ superfield for the ground state antitriplet baryons is simply identical to T_i. The superfields transform under parity as

$$\begin{aligned} \mathcal{R}_\mu(\vec{x},t) &\xrightarrow{P} \gamma_0 \mathcal{R}^\mu(-\vec{x},t) \\ \mathcal{S}_\mu(\vec{x},t) &\xrightarrow{P} -\gamma_0 \mathcal{S}^\mu(-\vec{x},t) \\ \mathcal{T}(\vec{x},t) &\xrightarrow{P} \gamma_0 \mathcal{T}(-\vec{x},t) \end{aligned} \tag{2.4}$$

and obey the constraints

$$\begin{aligned} \frac{1+\not{v}}{2}\mathcal{R}_\mu &= \mathcal{R}_\mu \qquad & \frac{1+\not{v}}{2}\mathcal{S}_\mu &= \mathcal{S}_\mu \qquad & \frac{1+\not{v}}{2}\mathcal{T} &= \mathcal{T}. \\ v^\mu \mathcal{R}_\mu &= 0 & v^\mu \mathcal{S}_\mu &= 0 & & \end{aligned} \tag{2.5}$$

These conditions ensure that $\mathcal{R}_\mu$ and $\mathcal{S}_\mu$ contain six degrees of freedom while $\mathcal{T}$ has two. The degree of freedom count thus agrees with the number of states that the superfields represent [16].

The constraints in (2.5) also fix the shifts in the baryon superfields induced by the reparameterization transformation

$$\begin{aligned} v &\to v + \epsilon/M \\ k &\to k - \epsilon \end{aligned} \tag{2.6}$$

where $v \cdot \epsilon = 0$. This change of variables leaves invariant the total four-momentum $p = Mv + k$ of a heavy hadron and induces only an $O(1/M^2)$ correction to $v^2 = 1$. The method for determining the induced shifts in the baryon superfields is entirely analogous to that for their meson counterparts which has previously been discussed in ref. [12]. So we only quote the results here:

$$\begin{aligned} \delta\mathcal{R}_\mu &= \frac{\not{\epsilon}}{2M}\mathcal{R}_\mu - \frac{\epsilon^\nu \mathcal{R}_\nu}{M} v_\mu \\ \delta\mathcal{S}_\mu &= \frac{\not{\epsilon}}{2M}\mathcal{S}_\mu - \frac{\epsilon^\nu \mathcal{S}_\nu}{M} v_\mu \\ \delta\mathcal{T} &= \frac{\not{\epsilon}}{2M}\mathcal{T}. \end{aligned} \tag{2.7}$$

The requirement that the effective theory remain invariant under the transformations in (2.6) and (2.7) forbids certain terms from appearing in the chiral Lagrangian [15].

The heavy baryons in the $\mathcal{R}_\mu$, $\mathcal{S}_\mu$ and $\mathcal{T}$ multiplets can interact with one another via emission and absorption of light Goldstone bosons. The Goldstone bosons result from the spontaneous breaking of $SU(3)_L \times SU(3)_R$ chiral symmetry down to its diagonal $SU(3)_{L+R}$ flavor subgroup and appear in the pion octet

$$\boldsymbol{\pi} = \sum_{a=1}^{8} \pi^a T^a = \frac{1}{\sqrt{2}} \begin{pmatrix} \sqrt{\frac{1}{2}}\pi^0 + \sqrt{\frac{1}{6}}\eta & \pi^+ & K^+ \\ \pi^- & -\sqrt{\frac{1}{2}}\pi^0 + \sqrt{\frac{1}{6}}\eta & K^0 \\ K^- & \overline{K}^0 & -\sqrt{\frac{2}{3}}\eta \end{pmatrix}. \tag{2.8}$$

It is convenient to arrange these fields into the exponentiated matrix functions $\Sigma = e^{2i\boldsymbol{\pi}/f}$ and $\xi = e^{i\boldsymbol{\pi}/f}$ where the parameter f equals the pion decay constant $f_\pi = 93$ MeV at lowest order. The matrix functions transform under the chiral symmetry group as

$$\begin{aligned} \Sigma &\to L\Sigma R^\dagger \\ \xi &\to L\xi U^\dagger(x) = U(x)\xi R^\dagger \end{aligned} \tag{2.9}$$

where L and R represent global elements of $SU(3)_L$ and $SU(3)_R$ while $U(x)$ acts like a local $SU(3)_{L+R}$ transformation. We further define the vector and axial vector fields

$$\begin{aligned} \mathbf{V}^\mu &= \frac{1}{2}(\xi^\dagger\partial^\mu\xi + \xi\partial^\mu\xi^\dagger) = \frac{1}{2f^2}[\boldsymbol{\pi}, \partial^\mu\boldsymbol{\pi}] - \frac{1}{24f^4}\Big[\boldsymbol{\pi}, \big[\boldsymbol{\pi}, [\boldsymbol{\pi}, \partial^\mu\boldsymbol{\pi}]\big]\Big] + O(\boldsymbol{\pi}^6) \\ \mathbf{A}^\mu &= \frac{i}{2}(\xi^\dagger\partial^\mu\xi - \xi\partial^\mu\xi^\dagger) = -\frac{1}{f}\partial^\mu\boldsymbol{\pi} + \frac{1}{6f^3}\big[\boldsymbol{\pi}, [\boldsymbol{\pi}, \partial^\mu\boldsymbol{\pi}]\big] + O(\boldsymbol{\pi}^5) \end{aligned} \tag{2.10}$$

which transform inhomogeneously and homogeneously under $SU(3)_{L+R}$ respectively:

$$\begin{aligned} \mathbf{V}^\mu &\to U\mathbf{V}^\mu U^\dagger + U\partial^\mu U^\dagger \\ \mathbf{A}^\mu &\to U\mathbf{A}^\mu U^\dagger. \end{aligned} \tag{2.11}$$

The pions in (2.8) derivatively couple to the baryon matter fields via these vector and axial vector combinations.

It is straightforward to construct the lowest order effective Lagrangian which describes the low energy interactions between the Qqq baryons and Goldstone bosons. One simply

writes down all possible terms that are Lorentz invariant, light chiral and heavy quark spin symmetric, and parity even:

$$
\begin{aligned}
\mathcal{L}_v^{(0)} = \sum_{Q=c,b} \Big\{ & \overline{\mathcal{R}}^i_\mu \big(-iv\cdot\mathcal{D} + \Delta M_{\mathcal{R}}\big)\mathcal{R}^\mu_i + \overline{\mathcal{S}}^\mu_{ij}\big(-iv\cdot\mathcal{D} + \Delta M_{\mathcal{S}}\big)S^{ij}_\mu + \overline{\mathcal{T}}^i iv\cdot\mathcal{D}\,\mathcal{T}_i \\
& + ig_1\,\varepsilon_{\mu\nu\sigma\lambda}\overline{\mathcal{S}}^\mu_{ik}\,v^\nu (A^\sigma)^i_j(\mathcal{S}^\lambda)^{jk} + ig_2\,\varepsilon_{\mu\nu\sigma\lambda}\overline{\mathcal{R}}^{\mu i}\,v^\nu(A^\sigma)^j_i(\mathcal{R}^\lambda)_j \\
& + h_1\Big[\epsilon_{ijk}\overline{\mathcal{T}}^i(A^\mu)^j_l\mathcal{S}^{kl}_\mu + \epsilon^{ijk}\overline{\mathcal{S}}^\mu_{kl}(A_\mu)^l_j\mathcal{T}_i\Big] \\
& + h_2\Big[\epsilon_{ijk}\overline{\mathcal{R}}^{\mu i}\,v\cdot A^j_l\mathcal{S}^{kl}_\mu + \epsilon^{ijk}\overline{\mathcal{S}}^\mu_{kl}\,v\cdot A^l_j\mathcal{R}_{\mu i}\Big]\Big\}.
\end{aligned}
\tag{2.12}
$$

A few points about this zeroth order Lagrangian should be noted. Firstly, the common mass splitting between the excited and ground state antitriplet multiplets is absorbed into the parameter $\Delta M_{\mathcal{R}} = M_{\mathcal{R}} - M_{\mathcal{T}}$. Similarly, $\Delta M_{\mathcal{S}} = M_{\mathcal{S}} - M_{\mathcal{T}}$ represents the splitting between the ground state sextet and antitriplet multiplets. These parameters do not vanish in the zero or infinite heavy quark mass limits and therefore appropriately reside within the leading order chiral Lagrangian. Secondly, the coupling constants $g_{1,2}$ and $h_{1,2}$ in (2.12) are expected to be of order unity on general dimensional analysis grounds [17]. However, their precise numerical values are *a priori* unknown and must be fitted to experiment. Finally, we observe that there are no terms in (2.12) which mediate the single Goldstone boson transitions $R^{(*)} \to T\boldsymbol{\pi}$ and $T \to T\boldsymbol{\pi}$. Such processes violate heavy quark spin symmetry and occur only at next-to-leading order in the $1/m_Q$ expansion.

The current experimental status of the baryons appearing in the heavy hadron chiral Lagrangian is very uneven. Data on strange charmed baryons is in short supply, and several have not yet been discovered. In contrast, a number of experiments within the past year have filled in most of the nonstrange members of the antitriplet and sextet multiplets. We will therefore focus upon the zero strangeness baryons in the remainder of this work.

The energy levels of the observed $\Lambda_c^{(*)}$ and $\Sigma_c^{(*)}$ states in $\mathcal{R}_\mu$, $\mathcal{S}_\mu$ and $\mathcal{T}$ are illustrated in fig. 1. As indicated in the figure, we interpret the two recently observed excited charmed baryons as the $I = 0$ members of the $\mathcal{R}_\mu$ multiplet. In the absence of well-established names for these baryons, we adopt the nomenclature convention of ref. [18] and denote the $J^P = \frac{1}{2}^-$ and $J^P = \frac{3}{2}^-$ states as Λ_{c1} and Λ^*_{c1} respectively. Averaging over the ARGUS, CLEO and E687 values for their masses, we find that they lie 308.0 ± 2.0 MeV and 341.4 ± 0.4 MeV above Λ_c. The splitting between these two P-wave baryon masses is comparable in magnitude to that between their P-wave meson analogues $D_1(2421$ MeV$)$ and $D_2(2465$ MeV$)$. We will keep track of this phenomenologically important mass difference even though it represents an $O(1/m_c)$ effect.

The splitting between Σ_c^* and Λ_c displayed in fig. 1 comes from another recent experimental result. The SKAT group claims to have observed the $J^P = \frac{3}{2}^+$ Σ_c^{*++} baryon for the first time in their bubble chamber experiment which uses a broad-band neutrino beam [19]. While their mass finding $M_{\Sigma_c^*} = 2530 \pm 7$ MeV must be treated with caution until independently confirmed by another group, we will adopt their reasonable central value in our subsequent analysis. Fortunately, none of our results will sensitively depend upon the precise numerical value for the Σ_c^* mass.

Having set up the necessary machinery for studying the two new Λ_c^* baryons, we proceed to examine their strong and radiative decay modes in the following two sections.

3. Strong Decays of Λ_c^*

The strong decays of the newly discovered excited charmed baryons are well-suited for Chiral Perturbation Theory analysis. The relatively small masses of Λ_{c1} and Λ_{c1}^* above Λ_c kinematically restrict their strong decays to soft pion emission. We therefore expect the chiral Lagrangian derivative expansion to be well-behaved for these new particles. Moreover, isospin conservation forbids single pion transitions between $\Lambda_{c1}^{(*)}$ and Λ_c. The excited $I = 0$ baryons must instead decay via an intermediate $I = 1$ state down to the $I = 0$ ground state. The released energy $M_{\Lambda_{c1}^{(*)}} - M_{\Lambda_c}$ is thus shared by two pions. [1]

Angular momentum and parity considerations require single pion transitions between the $\mathcal{R}_\mu$ and $\mathcal{S}_\mu$ multiplets to go through $L = 0$ or $L = 2$ partial waves. The D-wave coupling arises from dimension-five operators in the next-to-leading order chiral Lagrangian whose effects are quite suppressed. The S-wave coupling on the other hand is implemented by the dimension-four term proportional to h_2 in (2.12) which links Λ_{c1} with Σ_c and Λ_{c1}^* with Σ_c^*. The h_2 operator consequently mediates the barely allowed transition $\Lambda_{c1} \to \Sigma_c\pi$ at the rate

$$\Gamma(\Lambda_{c1} \to \Sigma_c\pi) = \frac{h_2^2}{4\pi f^2} \frac{M_{\Sigma_c}}{M_{\Lambda_{c1}}} (M_{\Lambda_{c1}} - M_{\Sigma_c})^2 \sqrt{(M_{\Lambda_{c1}} - M_{\Sigma_c})^2 - m_\pi^2}. \tag{3.1}$$

This process occurs so close to threshold that small isospin violating mass differences between members of the pion and charmed Sigma baryon multiplets cannot be ignored in

[1] The analogous kinematics for excited P-wave mesons is much less favorable. For example, the splitting between the D_2 and D mesons is almost 600 MeV, and single pion transitions between these two states are allowed. The validity of lowest order Chiral Perturbation Theory in this case is dubious at best.

the phase space factors of (3.1). Using the values $M_{\Sigma_c^0} = 2452.0$ MeV, $M_{\Sigma_c^+} = 2453.4$ MeV and $M_{\Sigma_c^{++}} = 2453.1$ MeV [20], we find the partial widths

$$\begin{aligned} \Gamma(\Lambda_{c1}^+ \to \Sigma_c^0 \pi^+) &= 3.3 h_2^2 \text{ MeV} && (3.2a) \\ \Gamma(\Lambda_{c1}^+ \to \Sigma_c^+ \pi^0) &= 6.0 h_2^2 \text{ MeV} && (3.2b) \\ \Gamma(\Lambda_{c1}^+ \to \Sigma_c^{++} \pi^-) &= 1.4 h_2^2 \text{ MeV}. && (3.2c) \end{aligned}$$

The analogous single pion transitions between the $J = \frac{3}{2}$ baryons in $\mathcal{R}_\mu$ and $\mathcal{S}_\mu$ are kinematically forbidden.

Double pion decays of Λ_{c1} and Λ_{c1}^* down to the Λ_c ground state proceed at leading order via the two pole graphs displayed in fig. 2. In order to obtain convergent decay rates from these diagrams, we must take into account the nonzero widths

$$\begin{aligned} \Gamma_{\Sigma_c} &= \frac{h_1^2}{12\pi f^2} \frac{M_{\Lambda_c}}{M_{\Sigma_c}} \big[(M_{\Sigma_c} - M_{\Lambda_c})^2 - m_\pi^2 \big]^{3/2} \simeq 2.5 h_1^2 \text{ MeV} \\ \Gamma_{\Sigma_c^*} &= \frac{h_1^2}{12\pi f^2} \frac{M_{\Lambda_c}}{M_{\Sigma_c^*}} \big[(M_{\Sigma_c^*} - M_{\Lambda_c})^2 - m_\pi^2 \big]^{3/2} \simeq 24 h_1^2 \text{ MeV} \end{aligned} \qquad (3.3)$$

of the intermediate Σ_c and Σ_c^* resonances. Their propagators thus appear as

$$\begin{aligned} D_{\Sigma_c} &= \frac{i}{v \cdot k - (M_{\Sigma_c} - M_{\Lambda_c}) + i\Gamma_{\Sigma_c}/2} \Lambda_+ \\ D_{\Sigma_c^*}^{\mu\nu} &= \frac{i}{v \cdot k - (M_{\Sigma_c^*} - M_{\Lambda_c}) + i\Gamma_{\Sigma_c^*}/2} \Lambda_+^{\mu\nu} \end{aligned} \qquad (3.4)$$

where $\Lambda_+ = (1 + \not{v})/2$ and $\Lambda_+^{\mu\nu} = \big[-g^{\mu\nu} + v^\mu v^\nu + \frac{1}{3}(\gamma^\mu + v^\mu)(\gamma^\nu - v^\nu) \big] \Lambda_+$ denote spin-$\frac{1}{2}$ and spin-$\frac{3}{2}$ projection operators respectively. We must also include a symmetry factor of 1/2 in the angular integration over the pions' momenta to avoid double counting the two identical bosons in the final state. A straightforward computation then yields the dimensionless differential decay rate

$$\begin{aligned} \frac{d\Gamma\big(\Lambda_{c1}^{(*)} \to \Lambda_c \pi^a \pi^b\big)}{dE_1} &= \frac{\delta^{ab}}{192\pi^3} \Big(\frac{h_1 h_2}{f^2} \Big)^2 \frac{M_{\Lambda_c}}{M_{\Lambda_{c1}^{(*)}}} \sqrt{(E_1^2 - m_\pi^2)(E_2^2 - m_\pi^2)} \\ &\times \Bigg[\frac{E_1^2 (E_2^2 - m_\pi^2)}{\big(M_{\Lambda_{c1}^{(*)}} - M_{\Sigma_c^{(*)}} - E_1 \big)^2 + \Gamma^2_{\Sigma_c^{(*)}}/4} + \frac{(E_1^2 - m_\pi^2) E_2^2}{\big(M_{\Lambda_{c1}^{(*)}} - M_{\Sigma_c^{(*)}} - E_2 \big)^2 + \Gamma^2_{\Sigma_c^{(*)}}/4} \Bigg] \end{aligned} \qquad (3.5)$$

expressed in terms of the two pion energies E_1 and $E_2 = M_{\Lambda_{c1}^{(*)}} - M_{\Lambda_c} - E_1$ measured in the decaying body's rest frame. [2] Integrating over E_1, we obtain the total rate

$$\Gamma\big(\Lambda_{c1}^{(*)} \to \Lambda_c \pi^a \pi^b\big) = \frac{h_2^2 \delta^{ab}}{8\pi^2 f^2} \frac{M_{\Sigma_c^{(*)}}}{M_{\Lambda_{c1}^{(*)}}} \frac{I}{\Big[(M_{\Sigma_c^{(*)}} - M_{\Lambda_c})^2 - m_\pi^2\Big]^{3/2}} \tag{3.6}$$

where

$$\begin{aligned} I = \frac{\Gamma_{\Sigma_c^{(*)}}}{2} \int_{m_\pi}^{M_{\Lambda_{c1}^{(*)}} - M_{\Lambda_c} - m_\pi} & dE_1 \sqrt{(E_1^2 - m_\pi^2)(E_2^2 - m_\pi^2)} \\ \times \Bigg[& \frac{E_1^2 (E_2^2 - m_\pi^2)}{\Big(M_{\Lambda_{c1}^{(*)}} - M_{\Sigma_c^{(*)}} - E_1\Big)^2 + \Gamma^2_{\Sigma_c^{(*)}}/4} + \frac{(E_1^2 - m_\pi^2) E_2^2}{\Big(M_{\Lambda_{c1}^{(*)}} - M_{\Sigma_c^{(*)}} - E_2\Big)^2 + \Gamma^2_{\Sigma_c^{(*)}}/4} \Bigg]. \end{aligned} \tag{3.7}$$

Since we do not know the values of h_1 and h_2, we cannot extract precise quantitative predictions from eqns. (3.5) - (3.7). However, these formulas do provide useful qualitative insight into the P-wave baryons' strong decays. In fig. 3, we plot $h_2^{-2} d\Gamma(\Lambda_{c1}^{(*)} \to \Lambda_c \pi^0 \pi^0)/dE_1$ versus E_1 with h_1 set equal to unity. As can clearly be seen in the figure, $\Lambda_{c1} \to \Lambda_c \pi^0 \pi^0$ is dominated by the pole regions where the intermediate Σ_c^+ state is very close to being on-shell. Its integrated rate is thus well approximated by the single π^0 partial width in $(3.2b)$. The rate in the charged pion channel is similarly well approximated by the sum of the widths in $(3.2a)$ and $(3.2c)$. Indeed, evaluating the phase space integral in (3.7) using the narrow width approximation

$$\frac{\Gamma_{\Sigma_c}/2}{(M_{\Lambda_{c1}} - M_{\Sigma_c} - E)^2 + \big(\Gamma_{\Sigma_c}/2\big)^2} \simeq \pi \delta(M_{\Lambda_{c1}} - M_{\Sigma_c} - E), \tag{3.8}$$

we simply recover eqn. (3.1) for $\Gamma(\Lambda_{c1} \to \Sigma_c \pi)$ which is independent of coupling constant h_1.

Nonresonant contributions generate a slight dependence of $\Gamma(\Lambda_{c1} \to \Lambda_c \pi^a \pi^b)$ upon h_1 as shown in fig. 4. But the decay of the $J^P = \frac{1}{2}^-$ state may essentially be viewed as the two step process $\Lambda_{c1} \to \Sigma_c \pi$ followed by $\Sigma_c \to \Lambda_c \pi$. In contrast, the double pion decay of Λ_{c1}^* cannot be regarded as a sequential transition. The virtual Σ_c^* intermediate state

[2] In the infinite charm mass limit, the recoiling Λ_c baryon carries off momentum but no kinetic energy. The two pions thus share all of the energy released by the decaying $\Lambda_{c1}^{(*)}$. This situation is similar to bouncing a ball off the earth. The earth must recoil to conserve momentum, but the ball bounces back with practically all its original kinetic energy [21].

is very much off-shell and produces no large resonant contribution to $\Lambda_{c1}^* \to \Lambda_c \pi\pi$. As a result, the strong interaction partial width of the $J^P = \frac{3}{2}^-$ state is more than an order of magnitude smaller than that of its $J^P = \frac{1}{2}^-$ partner.

As advertised in the Introduction, these qualitative findings on the excited charm baryon decay modes are in basic accord with the recent CLEO results reported in refs. [2] and [4]. They thus bolster one's confidence in the interpretation of the two new states as Λ_c^* baryons. To make further progress however, we need width information to pin down the values of the coupling constants in chiral Lagrangian (2.12). ARGUS has set a 90% CL upper bound of 3.2 MeV on the width of Λ_{c1}^* [1]. Unfortunately, this limit places only a weak constraint on the allowed parameter space in the h_1-h_2 plane. As fig. 4 demonstrates, the true natural width of Λ_{c1}^* is most likely too narrow to be resolved by current experimental detectors. On the other hand, there is a much better chance that the Λ_{c1} resonance is wide enough to be measured. In the $I = 1$ sector, the width resolving prospects for the $J^P = \frac{1}{2}^+$ and $J^P = \frac{3}{2}^+$ members of the $\mathcal{S}_\mu$ doublet are just opposite those for the $J^P = \frac{1}{2}^-$ and $J^P = \frac{3}{2}^-$ members of $\mathcal{R}_\mu$. We are therefore hopeful that experimentalists will be able to fix some of the free parameters in the heavy baryon chiral Lagrangian in the near future.

4. Electromagnetic Decays of Λ_c^*

The only decay modes of the two new $\Lambda_{c1}^{(*)}$ baryons that have so far been experimentally observed are their double pion transitions to Λ_c. But as shown in fig. 1, these P-wave hadrons can also de-excite down to the ground state via single photon emission. Unlike the strong interaction processes, the radiative channels are not severely phase space suppressed. Moreover, they produce two rather than three bodies in the final state. So the inherently weaker strength of the electromagnetic transitions could be offset by their more favorable kinematics. We explore such a possibility in this section.

Electromagnetic interactions may be incorporated into Heavy Hadron Chiral Perturbation Theory by gauging a $U(1)_{EM}$ subgroup of the global $SU(3)_L \times SU(3)_R$ chiral symmetry group. All derivatives appearing in the velocity dependent effective Lagrangian are then promoted to covariant derivatives with respect to electromagnetism. The leading dimension-four operators in (2.12) cannot contribute to S-matrix elements between states containing real photons. So to study heavy meson and baryon radiative transitions, one

must include a number of dimension-five operators into the chiral Lagrangian. In refs. [22–24], the $M1$ transitions between ground state hadrons containing a single heavy quark were analyzed. We now extend this earlier work to investigate the $E1$ decays of the new $\Lambda_{c1}^{(*)}$ baryons.

In the low energy theory, short wavelength photons with energies greater than the chiral symmetry breaking scale are integrated out and only long wavelength modes are retained. At lowest order in the $1/m_Q$ expansion, photons couple to just the light brown muck inside Qqq baryons leaving the spins of their heavy quark constituents unaltered. Such couplings generate the following contributions to the effective Lagrangian:

$$\begin{aligned}\mathcal{L}_v^{EM} = \sum_{Q=c,b} \frac{e(\Lambda_\chi)}{\Lambda_\chi} \Big\{ & i c_R \overline{\mathcal{R}}^j_\mu \mathcal{Q}^i_j \mathcal{R}_{\nu i} F^{\mu\nu} + i c_S \overline{\mathcal{S}}_{\nu ij} \big(\mathcal{Q}^i_k \mathcal{S}^{kj}_\mu + \mathcal{Q}^j_k \mathcal{S}^{ik}_\mu \big) F^{\mu\nu} \\ & + c_{RS} \Big[\epsilon_{ijk} \overline{\mathcal{R}}^i_\mu \mathcal{Q}^j_l \mathcal{S}^{kl}_\nu + \epsilon^{ijk} \overline{\mathcal{S}}_{\nu,kl} \mathcal{Q}^l_j \mathcal{R}_{\mu i} \Big] \tilde{F}^{\mu\nu} + c_{RT} \Big[\overline{\mathcal{T}}^j \mathcal{Q}^i_j \mathcal{R}^\mu_i + \overline{\mathcal{R}}^{\mu i} \mathcal{Q}^j_i \mathcal{T}_j \Big] v^\nu F_{\mu\nu} \\ & + c_{ST} \Big[\epsilon_{ijk} \overline{\mathcal{T}}^i \mathcal{Q}^j_l \mathcal{S}^{kl}_\nu + \epsilon^{ijk} \overline{\mathcal{S}}_{\nu,kl} \mathcal{Q}^l_j \mathcal{T}_i \Big] v_\mu \tilde{F}^{\mu\nu} \Big\}. \end{aligned} \tag{4.1}$$

Here $F^{\mu\nu}$ and $\tilde{F}^{\mu\nu}$ are the electromagnetic field strength tensor and its dual, and $\boldsymbol{\mathcal{Q}} = \frac{1}{2}(\xi Q_{EM} \xi^\dagger + \xi^\dagger Q_{EM} \xi)$ where

$$Q_{EM} = \begin{pmatrix} \mathcal{Q}_u & & \\ & \mathcal{Q}_d & \\ & & \mathcal{Q}_s \end{pmatrix} = \begin{pmatrix} \frac{2}{3} & & \\ & -\frac{1}{3} & \\ & & -\frac{1}{3} \end{pmatrix} \tag{4.2}$$

denotes the light quark electric charge matrix. The transformation rule $\boldsymbol{\mathcal{Q}} \to U \boldsymbol{\mathcal{Q}} U^\dagger$ for this spurion field renders the terms in (4.1) chiral symmetric [25].

The interactions between low energy photons and light brown muck take place at a long distance scale Λ_χ. In Chiral Perturbation Theory, this parameter represents the chiral symmetry breaking scale whose numerical value is approximately 1000 MeV. The phenomenologically successful Nonrelativistic Quark Model suggests however that a more appropriate value for Λ_χ in $\mathcal{L}_v^{EM}$ is a typical constituent quark mass of 300 MeV. We cannot really distinguish between these two numbers at the level of naive dimensional analysis as they differ by only a factor of three. But since Λ_χ enters quadratically into radiative decay rates, it is important to minimize its uncertainty as much as possible. So we will compromise and take Λ_χ to be the geometric mean between the CPT and NRQM values:

$${\Lambda_\chi}^2 = (300 \text{ MeV})(1000 \text{ MeV}) = (547.7 \text{ MeV})^2. \tag{4.3}$$

Hopefully this guess for ${\Lambda_\chi}^2$ does not lie more than a factor of three away from the true value.

The terms in (4.1) proportional to c_{R}, c_{S} and c_{ST} mediate the $M1$ radiative transitions $R^* \to R\gamma$, $S^* \to S\gamma$ and $S^{(*)} \to T\gamma$. The c_{RS} and c_{RT} operators on the other hand generate $E1$ decays $R^{(*)} \to S^{(*)}\gamma$ and $R^{(*)} \to T\gamma$. After extracting the $\Lambda_{c1}^{(*)}$ components from these interactions, we find the following radiative partial widths:

$$\Gamma(\Lambda_{c1}^* \to \Lambda_{c1}\gamma) = \frac{4c_{R}^2}{81}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Lambda_{c1}}}{M_{\Lambda_{c1}^*}}\left(\frac{M_{\Lambda_{c1}^*}^2 - M_{\Lambda_{c1}}^2}{2M_{\Lambda_{c1}^*}}\right)^3 = 4.36 \times 10^{-5}\, c_{R}^2 \text{ MeV} \tag{4.4}$$

$$\Gamma(\Lambda_{c1} \to \Sigma_c\gamma) = \frac{8c_{RS}^2}{9}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Sigma_c}}{M_{\Lambda_{c1}}}\left(\frac{M_{\Lambda_{c1}}^2 - M_{\Sigma_c}^2}{2M_{\Lambda_{c1}}}\right)^3 = 0.052\, c_{RS}^2 \text{ MeV} \tag{4.5a}$$

$$\Gamma(\Lambda_{c1} \to \Sigma_c^*\gamma) = \frac{4c_{RS}^2}{9}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Sigma_c^*}}{M_{\Lambda_{c1}}}\left(\frac{M_{\Lambda_{c1}}^2 - M_{\Sigma_c^*}^2}{2M_{\Lambda_{c1}}}\right)^3 = 0.003\, c_{RS}^2 \text{ MeV} \tag{4.5b}$$

$$\Gamma(\Lambda_{c1}^* \to \Sigma_c\gamma) = \frac{2c_{RS}^2}{9}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Sigma_c}}{M_{\Lambda_{c1}^*}}\left(\frac{M_{\Lambda_{c1}^*}^2 - M_{\Sigma_c}^2}{2M_{\Lambda_{c1}^*}}\right)^3 = 0.024\, c_{RS}^2 \text{ MeV} \tag{4.5c}$$

$$\Gamma(\Lambda_{c1}^* \to \Sigma_c^*\gamma) = \frac{10c_{RS}^2}{9}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Sigma_c^*}}{M_{\Lambda_{c1}^*}}\left(\frac{M_{\Lambda_{c1}^*}^2 - M_{\Sigma_c^*}^2}{2M_{\Lambda_{c1}^*}}\right)^3 = 0.023\, c_{RS}^2 \text{ MeV} \tag{4.5d}$$

$$\Gamma(\Lambda_{c1} \to \Lambda_c\gamma) = \frac{4c_{RT}^2}{27}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Lambda_c}}{M_{\Lambda_{c1}}}\left(\frac{M_{\Lambda_{c1}}^2 - M_{\Lambda_c}^2}{2M_{\Lambda_{c1}}}\right)^3 = 0.103\, c_{RT}^2 \text{ MeV} \tag{4.6a}$$

$$\Gamma(\Lambda_{c1}^* \to \Lambda_c\gamma) = \frac{4c_{RT}^2}{27}\frac{\alpha_{EM}(\Lambda_\chi)}{{\Lambda_\chi}^2}\frac{M_{\Lambda_c}}{M_{\Lambda_{c1}^*}}\left(\frac{M_{\Lambda_{c1}^*}^2 - M_{\Lambda_c}^2}{2M_{\Lambda_{c1}^*}}\right)^3 = 0.078\, c_{RT}^2 \text{ MeV}. \tag{4.6b}$$

As required heavy quark spin symmetry, the sum of the widths in ($4.5a, b$) equals the sum of those in ($4.5c, d$) in the infinite charm mass limit [26]. Similarly, the rates in ($4.6a$) and ($4.6b$) become degenerate when $m_c \to \infty$.

The leading order results in eqns. (4.4) - (4.6) cannot be trusted to provide much more than order of magnitude estimates for the $\Lambda_{c1}^{(*)}$ radiative decay rates. Yet comparing these electromagnetic partial widths with their strong interaction counterparts, we can draw some general qualitative conclusions. Firstly, we expect on the basis of naive dimensional analysis that the unknown c_{R}, c_{RS} and c_{RT} couplings are of order unity. The numerical partial width estimates suggest that some of the electromagnetic branching fractions might be measurable. Referring to fig. 4, we see that the two pion decay mode of Λ_{c1} dominates

over its radiative channels. The electromagnetic branching fraction for the $J^P = \frac{1}{2}^-$ state is thus most likely less than a few percent. On the other hand, since the double pion width of $J^P = \frac{3}{2}^-$ Λ_{c1}^* is much more narrow, $\Gamma(\Lambda_{c1}^* \to \Lambda_c\gamma)/\Gamma(\Lambda_{c1}^* \to \Lambda_c\pi\pi)$ could be sizable and perhaps greater than unity. Finally, we note that the radiative mode $\Lambda_{c1}^* \to \Sigma_c^*\gamma$ may provide a means for detecting the Σ_c^* baryon. The branching fraction for this process is small but not negligible. A search for this transition could therefore yield evidence for the elusive $I = 1$, $J^P = \frac{3}{2}^+$ state.

5. Conclusion

The basic interpretation of the two new excited charm baryons as $I = 0$ members of an excited P-wave doublet holds together remarkably well. Since the splitting between $\Lambda_{c1}^{(*)}$ and Λ_c is relatively small, these excited hadrons are well suited for incorporation into Heavy Hadron Chiral Perturbation Theory. Many experimental and theoretical details clearly remain to be filled into the picture which we have outlined here. In particular, width and branching ratio information are needed to fix the several new parameters that enter into the excited baryon sector.

A number of extensions of this work could be pursued in the future. For example, the primary decay modes of the $\Xi_{c1}^{0(*)}$ and $\Xi_{c1}^{+(*)}$ partners of $\Lambda_{c1}^{(*)}$ ought to be analyzed. As we have seen, there is no leading order term which links any of the states in the $\mathcal{R}_\mu$ antitriplet superfield with members of the ground state $\mathcal{T}$ multiplet. So if kinematically allowed, single kaon decays of these P-wave strange charmed baryons down to Λ_c are suppressed by $1/m_c$. A theoretical study of the dominant $\Xi_{c1}^{(*)}$ transitions would help guide an experimental search for these states. Alternatively, one might consider including other excited P-wave Qqq baryons into the heavy chiral Lagrangian. There are many such states waiting to be discovered.

In short, excited heavy baryon physics is a subject in which we may look forward to experimental and theoretical progress in the near future.

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Figure Captions

Fig. 1. Lowest lying $I = 0$ and $I = 1$ charmed baryon states. Experimentally measured mass splittings in MeV of the baryons above the Λ_c(2286 MeV) ground state are indicated in parentheses. The dominant pion decay modes of the excited $J^P = \frac{1}{2}^-$ Λ_{c1} and $J^P = \frac{3}{2}^-$ Λ^*_{c1} states are illustrated by the solid and dashed lines respectively. Their allowed radiative transitions down to the ground state are represented by the squiggly curves.

Fig. 2. Leading order pole graphs that contribute to $\Lambda^{(*)}_{c1} \to \Lambda_c \pi\pi$.

Fig. 3. Dimensionless differential decay rates $h_2^{-2} d\Gamma(\Lambda_{c1} \to \Lambda_c \pi^0 \pi^0)/dE_1$ (solid curve) and $10 \times h_2^{-2} d\Gamma(\Lambda^*_{c1} \to \Lambda_c \pi^0 \pi^0)/dE_1$ (dashed curve) plotted against pion energy E_1 measured in the excited charm baryon's rest frame. The coupling constant h_1 is set equal to unity in this graph.

Fig. 4. Integrated double pion decay rates of Λ_{c1} (solid curve) and Λ^*_{c1} (dashed curve) plotted as functions of coupling constant h_1. The neutral and charged pion channel contributions are summed together in this graph.